Δs density in a proton and unpolarized lepton–polarized proton scatterings

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Abstract

It is shown that the parity–violating deep–inelastic scatterings of unpolarized charged leptons on polarized protons, $\ell^{\mp} + \vec{P} \rightarrow \overset{(-)}{\nu_{\ell}} + X$, could provide a sensitive test for the behavior and magnitude of the polarized strange–quark density in a proton. Below charm threshold these processes are also helpful to uniquely determine the magnitude of individual polarized parton distributions.

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There have been interesting problems on strange–quark (s–quark) contents in contemporary hadron physics. Deductions of the σ–term from pion-nucleon scatterings imply an existence of significant s–quark contents in a nucleon[1]. New analysis suggests that about one third of the rest mass of the proton comes from $s\bar{s}$ pairs. So far, interesting experimental proposals[2] have been presented to measure the neutral weak form factors of the nucleon which might be sensitive to the s–quarks inside the nucleon. A different idea is also proposed to directly probe the s–quark content of the proton by using the lepto– and photo–production of ϕ–meson that is essentially 100% $s\bar{s}$[3]. Another surprizing results on the s–quark contents in the nucleon have been drawn from the data of polarized deep inelastic scatterings[4]. To our surprise, the experimental data have suggested that, contrary to the prediction of the naive quark model, there is a large and negative contribution of s–quarks to the proton spin, *i.e.* $\Delta s = -0.12$, and furthermore very little of the proton spin is carried by quarks. For low-energy properties of baryons, conventional phenomenological quark models treat nucleons as consisting of only u– and d–quarks and thus it naturally comes as a big surprise when some recent measurements and theoretical analyses have indicated a possible existence of a sizable s–quark. In order to get deep understanding of hadron dynamics, it is very important to investigate the behavior of s–quarks in a nucleon. In this paper, we concentrate on the behavior of the polarized s–quark and study the processes sensitive to its polarized distributions in the nucleon.

So far, several people have suggested various processes, which are sensitive to polarized s–quark distributions, such as Drell–Yan processes[5], inclusive $W^{\pm}-$ and Z^0–productions[6] in polarized proton–polarized proton collisions, and also inclusive $\pi^{\pm}-$ and $K^{\pm}$–productions in polarized lepton–polarized proton scatterings[7]. However, since the differential cross sections for Drell–Yan processes and inclusive $W^{\pm}-/Z^0$–hadroproductions are described by the product of two parton distributions participating in such processes, one cannot extract the x–dependence of polarized s–quark distributions without ambiguities from such cross sections. In addition, those of inclusive $\pi^{\pm}-$ and $K^{\pm}$–leptoproductions include the fragmentation func-

tions of $\pi^{\pm}-$ and $K^{\pm}$–decays which possess some theoretical ambiguities, and hence it is also difficult to derive the exact behavior of polarized s–quark distributions from these processes. Recently, it has been pointed out that parity–violating polarized electron elastic scatterings on unpolarized protons can give informations on the matrix elements, $\langle p|\bar{s}\Gamma_\mu s|p\rangle$ with Γ_μ $= \gamma_\mu$ and $\gamma_\mu\gamma_5$[8]. However, since its differential cross section includes not only the spin–dependent but also spin–independent proton form factors, one cannot extract the polarized s–quark content without ambiguities even from such processes. Here we consider a different process for examining the polarized s–quark density, which is the parity–violating polarized deep inelastic scattering at high energy. It must be advantageous to study such a process because its differential cross section includes only the spin–dependent structure function of the proton and is explicitly described as a function of x.

In parity–violating deep inelastic scatterings of unpolarized charged lepton on longitudinally polarized proton, an interesting parameter is the single–spin asymmetry $A_L^{W\mp}$ defined as

$$
\begin{aligned}
A_L^{W\mp} &= \frac{(d\sigma_{++}^{W\mp} + d\sigma_{-+}^{W\mp}) - (d\sigma_{+-}^{W\mp} + d\sigma_{--}^{W\mp})}{(d\sigma_{++}^{W\mp} + d\sigma_{-+}^{W\mp}) + (d\sigma_{+-}^{W\mp} + d\sigma_{--}^{W\mp})} \\
&= \frac{d\sigma_{0+}^{W\mp} - d\sigma_{0-}^{W\mp}}{d\sigma_{0+}^{W\mp} + d\sigma_{0-}^{W\mp}} = \frac{d\Delta_L\sigma^{W\mp}/dx}{d\sigma^{W\mp}/dx} ,
\end{aligned}
\tag{1}
$$

where $d\sigma_{0-}^{W\mp}$, for instance, denotes that the lepton is unpolarized and the helicity of the proton is negative. Note that since a fast incoming negatively (positively) charged lepton, ℓ^- (ℓ^+), couples to a W–boson only when it has a negative (positive) helicity, part of spin–dependent cross sections in eq.(1) should be zero. For parity–violating weak–interacting reactions with $W^\mp$ exchanges, $\ell^\mp + \vec{P} \to \overset{(-)}{\nu_\ell} + X$, the spin–dependent and spin–independent differential cross sections as a function of momentun fraction x are given by[9]

$$
\frac{d\Delta_L\sigma^{W\mp}}{dx} = 16\pi M_N E \frac{\alpha^2}{Q^4}\eta\left\{\pm(\frac{2}{3} + \frac{xM_N}{6E})x\, g_1^{W\mp}(x, Q^2) + (\frac{2}{3} - \frac{xM_N}{12E})\, g_3^{W\mp}(x, Q^2)\right\} ,
\tag{2}
$$

$$
\frac{d\sigma^{W\mp}}{dx} = 16\pi M_N E \frac{\alpha^2}{Q^4}\eta\left\{(\frac{2}{3} - \frac{xM_N}{4E})\, F_2^{W\mp}(x, Q^2) \pm \frac{1}{3}x\, F_3^{W\mp}(x, Q^2)\right\} ,
\tag{3}
$$

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where E is the energy of the charged lepton beam and M_N the mass of the proton. η is written in terms of the W–boson mass M_W as

$$\eta = \frac{1}{2}\left(\frac{G_F M_W^2}{4\pi\alpha}\frac{Q^2}{Q^2 + M_W^2}\right) . \tag{4}$$

$g_1^{W^\mp}$, $g_3^{W^\mp}$ in eq.(2) and $F_2^{W^\mp}$, $F_3^{W^\mp}$ in eq.(3) represent spin–dependent and spin–independent proton structure functions, respectively. Below charm threshold, the region of which could be investigated by SMC and/or E143 Collabolations, we can describe these structure functions for the W^- exchange as

$$F_2^{W^-}(x, Q^2) = 2x\left[c_1\{u_v(x, Q^2) + u_s(x, Q^2)\} + c_2\ \bar{d}_s(x, Q^2) + c_3\ \bar{s}_s(x, Q^2)\right] ,$$

$$F_3^{W^-}(x, Q^2) = 2\left[c_1\{u_v(x, Q^2) + u_s(x, Q^2)\} - c_2\ \bar{d}_s(x, Q^2) - c_3\ \bar{s}_s(x, Q^2)\right] , \tag{5}$$

$$g_1^{W^-}(x, Q^2) = \left[c_1\{\delta u_v(x, Q^2) + \delta u_s(x, Q^2)\} + c_2\ \delta\bar{d}_s(x, Q^2) + c_3\ \delta\bar{s}_s(x, Q^2)\right] ,$$

$$g_3^{W^-}(x, Q^2) = 2x\left[c_1\{\delta u_v(x, Q^2) + \delta u_s(x, Q^2)\} - c_2\ \delta\bar{d}_s(x, Q^2) - c_3\ \delta\bar{s}_s(x, Q^2)\right] ,$$

and similarly for the W^+ exchange

$$F_2^{W^+}(x, Q^2) = 2x\left[c_1\ \bar{u}_s(x, Q^2) + c_2\{d_v(x, Q^2) + d_s(x, Q^2)\} + c_3\ s_s(x, Q^2)\right] ,$$

$$F_3^{W^+}(x, Q^2) = 2\left[-c_1\ \bar{u}_s(x, Q^2) + c_2\{d_v(x, Q^2) + d_s(x, Q^2)\} + c_3\ s_s(x, Q^2)\right] , \tag{6}$$

$$g_1^{W^+}(x, Q^2) = \left[c_1\ \delta\bar{u}_s(x, Q^2) + c_2\{\delta d_v(x, Q^2) + \delta d_s(x, Q^2)\} + c_3\ \delta s_s(x, Q^2)\right] ,$$

$$g_3^{W^+}(x, Q^2) = 2x\left[-c_1\ \delta\bar{u}_s(x, Q^2) + c_2\{\delta d_v(x, Q^2) + \delta d_s(x, Q^2)\} + c_3\ \delta s_s(x, Q^2)\right]$$

with CKM matrix elements

$$c_1 = |U_{ud}|^2 + |U_{us}|^2 , \quad c_2 = |U_{ud}|^2 , \quad c_3 = |U_{us}|^2 . \tag{7}$$

Here $\delta q(x, Q^2) = q_+(x, Q^2) - q_-(x, Q^2)$ $(q(x, Q^2) = q_+(x, Q^2) + q_-(x, Q^2))$ stands for the polarized (unpolarized) quark distribution in a proton, and $q_+(x, Q^2)$ $(q_-(x, Q^2))$ the quark density having a momentum fraction x with the helicity parallel (anti–parallel) to the proton helicity. It should be noticed that since $g_1^{W^-}$ and $g_1^{W^+}$ have no relation to the flavor singlet axial–vector current, it is not affected by the axial anomaly which, in deep inelastic reactions

with one–photon exchanges, leads to rather large contributions of the polarized gluons to polarized quark distribution functions[10].

In order to examine how the observed parameter is affected by the behavior of polarized s–quark distributions, we calculate $A_L^{W^\mp}$ by substituting eqs.(2) and (3) into eq.(1) as follows,

$$A_L^{W^\mp} = \frac{\pm(\frac{2}{3} + \frac{xM_N}{6E})x\ g_1^{W^\mp}(x,Q^2) + (\frac{2}{3} - \frac{xM_N}{12E})\ g_3^{W^\mp}(x,Q^2)}{(\frac{2}{3} - \frac{xM_N}{4E})\ F_2^{W^\mp}(x,Q^2) \pm \frac{1}{3}x\ F_3^{W^\mp}(x,Q^2)}\ . \qquad (8)$$

As typical examples of the polarized s–quark distribution functions, we take the following three different types; (i) negative and large Δs with $\Delta s(Q^2 = 4\text{GeV}^2) = -0.11$ (BBS model)[11] , (ii) zero Δs with $\Delta s(Q^2 = 10\text{GeV}^2)=0$[12] , (iii) positive and small Δs with $\Delta s(Q^2 = 10\text{GeV}^2) = 0.02$ (KMTY model)[13] , where $\Delta s(Q^2)$ is the first moment of $\delta s(x,Q^2)$ and its value means the amount of the proton spin carried by the s–quark. All of the models, (i), (ii) and (iii), can reproduce the EMC and SMC data on $xg_1^p(x)$ equally well. The x–dependence of these distributions are depicted at $Q^2 = 10\text{GeV}^2$ in fig.1. (The explicit forms of $\delta s(x,Q^2)$ are presented in respective references.) By using the polarized s–quark distribution of each type, together with the polarized and unpolarized parton distribution of BBS parametrization[11] for (i), of Cheng–Lai[12] and DFLM parametrizaton[14] for (ii), and of KMTY[13] and DO parametrization[15] for (iii), we estimate the $A_L^{W^\mp}$ at a typical charged lepton energy $E = 180\text{GeV}$ and momentum transfer squared $Q^2 = 10\text{GeV}^2$ whose kinematical region can be covered by SMC experiments. We see that $A_L^{W^-}$ for the W^- exchange process depends on the behavior of polarized s–quark distributons little because the spin–dependent proton structure functions, $g_1^{W^-}(x,Q^2)$ and $g_3^{W^-}(x,Q^2)$, included in $A_L^{W^-}$ are dominantly controlled by the polarized valence u–quark ditribution $\delta u_v(x,Q^2)$ whose magnitude Δu_v is much larger than $\Delta \bar{s}_s$ in the proton. However, the situation is quite different for $g_1^{W^+}(x,Q^2)$ and $g_3^{W^+}(x,Q^2)$ originated from W^+ exchanges because $g_1^{W^+}(x,Q^2)$ and $g_3^{W^+}(x,Q^2)$ have no contribution of the polarized valence u–quark distribution. Although $g_1^{W^+}(x,Q^2)$ and $g_3^{W^+}(x,Q^2)$ contain the polarized valence d–quark distribution $\delta d_v(x,Q^2)$, the absolute value of Δd_v is quite smaller than that of Δu_v. Thus, $A_L^{W^+}$ is expected to be rather sensitive to the polarized s–quark distribution function compared to the case of

$A_L^{W^-}$. The x–dependence of $xg_1^{W^+}$, $g_3^{W^+}$ and $A_L^{W^+}$ are calculated and shown in figs.2, 3 and 4, respectively. From fig.4, one can observe that the behavior of $A_L^{W^+}$ significantly depends on the polarized s–quark distribution for not very small x regions and hence we can distinguish the model of polarized s–quark distributions from the data of $A_L^{W^+}$. The reader might consider that the difference obtained here could be originated from our procedure of having used the different unpolarized parton distributions for respective models, (i), (ii) and (iii). In order to examine if the results are really meaningful, we have carried out the same calculation using in common the DFLM parametrization for unpolarized distributions as an example. Calculated results are presented by lines with small circles in fig.4. From this result, we can say that the conclusion remains to be unchanged.

However, it should be noted that one cannot directly extract the s–quark distribution from $xg_1^{W^+}(x, Q^2)$ and $g_3^{W^+}(x, Q^2)$. This is because, as shown in eq.(2), the differential cross sections are described by a linear combination of $xg_1^{W^+}(x, Q^2)$ and $g_3^{W^+}(x, Q^2)$ and we cannot measure independently $xg_1^{W^+}(x, Q^2)$ and $g_3^{W^+}(x, Q^2)$ in experiments. It is interesting to recall that the situation is quite different for the case of one–photon exchange processes. Although the differential cross sections are described by the linear combination of $g_1^p(x)$ and $g_2^p(x)$ in that case as well, $g_2^p(x)$ can be kinematically neglected compared to $g_1^p(x)$ and hence one can easily measure $g_1^p(x)$ in experiment. This is not the case for $W^\pm$ exchange processes. The information on the polarized s–quark distribution reflects to $A_L^{W^+}$ which contains $xg_1^{W^+}(x, Q^2)$ and $g_3^{W^+}(x, Q^2)$. In practice, the differential cross sections for these processes are small because of weakly interacting W–boson exchanges: for example, $d\Delta_L\sigma^{W^+}/dx = -0.093\ (-0.011, -0.099)$ [pb] and $d\sigma^{W^+}/dx = 0.62\ (0.75, 0.52)$ [pb] for the type of (iii) ((i), (ii)) at $x = 0.1$, $E = 180\text{GeV}$ and $Q^2 = 10\text{GeV}^2$. Therefore, from the experimental point of view, we must have high luminosity in order to get informations on the polarized s–quark distribution function $\delta s(x, Q^2)$ inside a proton. Futhermore it might be practically rather difficult to identify the missing events from $\ell^+ + \vec{P} \to \bar{\nu}_\ell + X$. However, we believe that these difficulties can be technically overcome.

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Another importance for getting $g_1^{W^\pm}$ comes from the fact that, below charm threshold, one can uniquely determine the magnitude of individual polarized parton density by combining the data on $g_1^{W^\pm}$ with the ones on neutron β–decay, hyperon β–decay and the spin–dependent structure function of proton g_1^p. The following combinations of individual polarized parton content are well–known,

$$\Delta u - \Delta d = a_3 \ , \tag{9}$$

$$\Delta u + \Delta d - 2\Delta s = a_8 \ , \tag{10}$$

$$\frac{4}{18}\Delta u + \frac{1}{18}\Delta d + \frac{1}{18}\Delta s - \frac{\alpha_s}{6\pi}\Delta g = a_0^\gamma \ , \tag{11}$$

where Δg in eq.(11) represents the amount of the proton spin carried by gluons and is introduced by taking account of the $U_A(1)$ anomaly of QCD[10]. Although the values of a_3, a_8 and a_0^γ are known from the experimental data on neutron β–decay, hyperon β–decay and the spin–dependent structure function of proton g_1^p, respectively, it is impossible to determine the magnitude of individual content uniquely from these equations alone since there exist four independent variables for three equations. However, if the values of the first moment of $g_1^{W^+}(x)$ and $g_1^{W^-}(x)$ can be obtained experimentally

$$\int_0^1 g_1^{W^+}(x)dx + \int_0^1 g_1^{W^-}(x)dx = c_1 \ \Delta u + c_2 \ \Delta d + c_3 \ \Delta s = a_0^{W^+} + a_0^{W^-} \ , \tag{12}$$

then, from four independent equations, (9), (10), (11) and (12), Δu, Δd, Δs and Δg can be determined as follows,

$$\Delta u = \frac{(2c_2 + c_3)a_3 + c_3 a_8 + 2(a_0^{W^+} + a_0^{W^-})}{2(c_1 + c_2 + c_3)} \ , \tag{13}$$

$$\Delta d = \frac{(-2c_1 - c_3)a_3 + c_3 a_8 + 2(a_0^{W^+} + a_0^{W^-})}{2(c_1 + c_2 + c_3)} \ , \tag{14}$$

$$\Delta s = \frac{(-c_1 + c_2)a_3 - (c_1 + c_2)a_8 + 2(a_0^{W^+} + a_0^{W^-})}{2(c_1 + c_2 + c_3)} \ , \tag{15}$$

$$\Delta g = -\frac{\pi}{3\alpha_s}\frac{1}{2(c_1 + c_2 + c_3)}\{3(c_1 - 3c_2 - c_3)a_3 \tag{16}$$
$$+(c_1 + c_2 - 5c_3)a_8 + 36(c_1 + c_2 + c_3)a_0^\gamma - 12(a_0^{W^+} + a_0^{W^-})\} \ .$$

But how can we actually measure $a_0^{W^\pm}$ in experiment? As long as we remain in the experiment with the longitudinally polarized proton target, we cannot determine $a_0^{W^\pm}$ experimentally. However, if we consider the cross section with the transversely polarized proton target, one can obtain $a_0^{W^\pm}$ as described in the following.

The formulas of the transversely polarized cross section have been given by Anselmino et al.[9] as follows,

$$
\begin{aligned}
\frac{d\Delta_T \sigma^{W^\pm}}{dxdyd\phi} &= 32 M_N \frac{\alpha^2}{Q^4} \eta \cos(\psi - \phi) \sqrt{xy M_N \{2(1-y)E - xy M_N\}} \\
&\times\ x(1-y)\left(\pm g_1^{W^\pm}(x, Q^2) + \frac{1}{2x}\, g_3^{W^\pm}(x, Q^2)\right)\ ,
\end{aligned}
\tag{17}
$$

where ϕ is the azimuthal angle of the lepton in the final state, and ψ an angle between the proton spin and x–axis in the xy plane orthogonal to the lepton direction (z–axis). These angles must be fixed in principle by suitably arranging experimental apparatus. $d\Delta_T \sigma^{W^\pm}/dxdyd\phi$ is defined as

$$
\frac{d\Delta_T \sigma^{W^\pm}}{dxdyd\phi} = \frac{d\sigma_{0\uparrow}^{W^\pm}}{dxdyd\phi} - \frac{d\sigma_{0\downarrow}^{W^\pm}}{dxdyd\phi}\ ,
\tag{18}
$$

where $d\sigma_{0\uparrow}^{W^\pm}$ ($d\sigma_{0\downarrow}^{W^\pm}$) denotes that the lepton is unpolarized and the proton is transversely polarized with its spin orthogonal to the lepton direction at an angle ψ ($\psi + \pi$) to the x–axis. By integrating eq.(17) over y, one can easily derive the following formula,

$$
\begin{aligned}
\frac{d\Delta_T \sigma^{W^+}}{dxd\phi} - \frac{d\Delta_T \sigma^{W^-}}{dxd\phi} &= C\ x^{3/2}\left[\int_0^1 dy\sqrt{y}(1-y)\sqrt{1-y-\frac{xy M_N}{2E}}\right] \\
&\times\ \left(g_1^{W^+}(x, Q^2) + g_1^{W^-}(x, Q^2)\right)\ ,
\end{aligned}
\tag{19}
$$

with $C = 32\sqrt{2E}M_N^{3/2}\frac{\alpha^2}{Q^4}\eta \cos(\psi - \phi)$. In eq.(19), the integral in the square bracket depends on x alone and can be written by $f(x)$. Practically, $f(x)$ can be very nicely approximated by the formula,

$$
f(x) = 0.19635\left(1 - \frac{2.45}{E}x\right)\ ,
\tag{20}
$$

which reproduces the exact result with accuracy better than 10^{-4} for $E > 50$GeV. Then, we can get

$$
g_1^{W^+}(x) + g_1^{W^-}(x) = \frac{d\Delta_T \sigma^{W^+}/dxd\phi - d\Delta_T \sigma^{W^-}/dxd\phi}{C\ x^{3/2}\ f(x)}\ ,
\tag{21}
$$

where the right–hand side of eq.(21) can be determined from experiment. Therefore, if we carry out the experiment with the transversely polarized target, we can obtain the sum of $a_0^{W^-}$ and $a_0^{W^+}$ by integrating eq.(21) over x, which leads to the unique determination of the magnitude of individual polarized parton densities in a proton. On the contrary, we can predict the values of $a_0^{W^-}$ and $a_0^{W^+}$ by using each type of polarized s–quark density. Some examples are given $a_0^{W^-} = -0.400$ and $a_0^{W^+} = 0.798$ for type (i), -0.282 and 0.956 for type (ii), -0.196 and 0.943 for type (iii) at $E = 180\text{GeV}$ and $Q^2 = 10\text{GeV}^2$. These predictions should be tested in the forthcoming experiment.

In summary, we have discussed the processes sensitive to the polarized s–quark distribution and have found that parity–violating reactions with W^+–boson exchange, $\ell^+ + \vec{P} \to \bar{\nu}_\ell + X$, are quite promising for giving us informations on polarized s–quark distribution functions inside a proton. Since the single–spin asymmetry $A_L^{W^+}$ for these processes significantly depends on the behavior of polarized s–quark distributions, one can test the behavior and magnitude of the s–quark polarization by measuring this quantity in experiments. Futhermore, we have shown that the amount of each quark and gluon carrying the proton spin can be uniquely determined if the spin–dependent proton structure functions, $g_1^{W^+}$ and $g_1^{W^-}$, are obtained below charm threshold by carrying out the experiment with the transversely polarized target.

Informations on polarized s–quark distributions are decisively important to understand the proton spin strucutre. We hope the present predictions could be tested in the forthcoming experiments.

Acknowledgements

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Figure captions

Fig. 1: The x–dependence of $x\delta s(x, Q^2)$ at $Q^2 = 10\text{GeV}^2$ for various types of polarized s–quark distributions (i)–(iii). (See text.) Type (i) is evolved up to $Q^2 = 10\text{GeV}^2$. The dash–dotted, dashed and solid lines indicate the results calculated for types (i), (ii) and (iii), respectively.

Fig. 2: The x–dependence of spin–dependent proton structure functions $xg_1^{W^+}$ for partiy–violating reactions with W^+–boson exchanges at $Q^2 = 10\text{GeV}^2$ for various types of Δs. Various lines represent the same as in fig.1.

Fig. 3: The spin–dependent proton structure functions $g_3^{W^+}$ as a function of x at $Q^2 = 10\text{GeV}^2$. Various lines represent the same as in fig.1.

Fig. 4: The single–spin asymmetry $A_L^{W^+}$ as a function of x at $E = 180\text{GeV}$ and $Q^2 = 10\text{GeV}^2$. The dash–dotted, dashed and solid lines correspond to types (i), (ii) and (iii), respectively. The lines added small circles for respective types are the results calculated by using in common the DFLM parametrization for unpolarized distributions.

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